

Connection between Dark Matter Abundance and Primordial Tensor Perturbations

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ABSTRACT

Primordial inflation and Dark Matter (DM) could both belong to the hidden sector. It is therefore plausible that the inflaton, which drives inflation, could couple to the DM either directly or indirectly, thus providing a common origin for both luminous and non-luminous matter. We explore this interesting possibility and show that, in certain scenarios, the DM mass can be correlated with the tensor-to-scalar ratio. This correlation might provide us with a window of opportunity for unravelling the properties of DM beyond the standard freeze-out paradigm.

Key Words: Inflation, Dark matter, Tensor modes.

1. INTRODUCTION

Primordial inflation is one of the simplest paradigms to explain the formation of large scale structures in our Universe and the observed features of the Cosmic Microwave Background spectrum (Ade et al., 2015b). Inflation is driven by the vacuum energy density of a scalar field, known as the inflaton, whose origin usually requires some beyond the Standard Model (SM) physics (Mazumdar and Rocher, 2011). On the other hand, various astrophysical and cosmological observations (Bertone et al., 2005) strongly suggest the existence of a non-luminous, non-baryonic form of matter, known as Dark Matter (DM). Although the masses and interactions of either inflaton or DM are still unknown, both must couple to the SM in some way, if not directly. Technically speaking, they can be considered to be SM gauge singlets, and therefore, could both belong to the dark or the hidden sector. In this case, one can imagine that the net DM, which we assume to be cosmologically stable, can be created via two processes:

(a) Decay: The inflaton ϕ could directly couple to the DM χ via renormalizable interactions, as shown in Fig. 1(a). For concreteness, we assume a fermionic DM, so that the interaction is of the Yukawa type, i.e. $\phi\bar{\chi}\chi$. We could also have a scalar DM with a trilinear coupling to inflaton. Apart from the inflaton itself, any heavy hidden sector scalar field X could also directly decay to DM via renormalizable interactions.

(b) Scattering: Inflaton must couple to the SM degrees of freedom (d.o.f) for the success of Big Bang Nucleosynthesis. Since the SM fermions are chiral, a SM singlet inflaton can in principle couple to the right-handed (RH) fermions ψ_R via renormalizable interactions of the form $\phi\bar{\psi}_R\psi_R$. Therefore, we can also create DM via *inflaton mediation*, as shown in Fig. 1(b). One can generalize this scenario to envisage that any heavy scalar mediator X could connect the SM d.o.f with the dark sector.

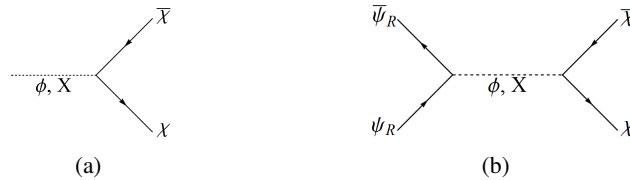


Figure 1: DM production from (a) non-thermal decay of inflaton/heavy scalar and (b) thermal scatterings with the SM d.o.f mediated by inflaton/heavy scalar field.

A rather natural outcome of this simple scenario is that the scale of inflation, determined by the inflaton potential $V(\phi)$, can be correlated with the DM properties in a rather intriguing way. If, for some

reason, the DM does not fully thermalize with the primordial plasma during its evolution, it can in principle retain the memory of how it was excited at the first instance, either (a) directly via the inflaton decay or (b) indirectly via scatterings mediated by the inflaton or a heavy scalar field. In case (a), the DM is essentially decoupled from the thermal bath since its creation. This leads to a non-thermal DM scenario, where the DM relic abundance is directly determined by the initial inflaton energy density (Allahverdi and Drees, 2002a,b). In case (b), if the effective coupling of the DM to the SM d.o.f is too small to fully thermalize the DM with the bath, but sufficient enough to produce the observed abundance of DM, this leads to the Feebly Interacting Massive Particle (FIMP) or freeze-in DM scenario (Hall et al., 2010). In both cases, the final DM relic density is sensitive to the initial conditions set by inflation (Dev et al., 2014), unlike in the standard thermal weakly interacting massive particle (WIMP) scenario (Kolb and Turner, 1990), thereby providing the unique possibility to directly link the DM properties with inflation.

In this paper, we show that in the non-WIMP scenarios, it is indeed possible to establish a connection between the DM and inflaton sectors via the scale of inflation, which can be determined by measuring the primordial tensor-to-scalar ratio. Note that for the WIMP DM scenario with a relatively large DM coupling to the SM d.o.f, there are many observational constraints from direct/indirect searches (Bertone et al., 2005), but few are applicable to the non-standard DM scenarios (a) and (b) discussed above. Therefore, the window of opportunity established in this paper is extremely useful for probing such DM candidates. We illustrate this for a simple class of inflationary models with an α -attractor potential (De Felice et al., 2011; Kallosh et al., 2013), but our results could be easily extended to other inflationary potentials.

2. A BRIEF REVIEW ON INFLATIONARY SET UP

Typically, the scale of inflation can be observationally determined by the tensor-to-scalar ratio $r = \mathcal{P}_T / \mathcal{P}_S$, where \mathcal{P}_S is the amplitude of the scalar perturbations given by $\mathcal{P}_S = 2.142_{-0.049}^{+0.049} \times 10^{-9}$ (Ade et al., 2015b) and \mathcal{P}_T is the amplitude of the tensor power spectrum. The Hubble rate $H = V(\phi)/3M_P^2$, where $M_P = 2.4 \times 10^{18}$ GeV is the reduced Planck mass, is given by (Kolb and Turner, 1990)

$$H \simeq 3 \times 10^{-5} \left(\frac{r}{0.1} \right)^{1/2} M_P. \quad (1)$$

Thus, it is possible that by measuring r one might get some insight about the DM relic abundance using the production mechanisms given in Fig. 1. In order to illustrate this, let us consider a class of inflationary models given by the following potential (De Felice et al., 2011; Kallosh et al., 2013):

$$V(\phi) = \frac{3}{4} M_P^2 M^2 \left(1 - e^{-\sqrt{\frac{2}{3\alpha}} \frac{\phi}{M_P}} \right)^2 \equiv \frac{3}{4} M_P^2 M^2 [1 - x(\phi)]^2, \quad (2)$$

where α is a free parameter and M is some mass scale governing inflation. For $\alpha = 1$, this is just the Starobinsky model (Starobinsky, 1980), whereas in the limit $\alpha \rightarrow \infty$, Eq. (2) reduces to the simple quadratic chaotic inflation potential (Linde, 1983) with constant mass, $M/\sqrt{\alpha}$. In this class of models, inflation occurs above the scale of M_P and terminates at

$$\phi_{\text{end}} \simeq \sqrt{\frac{3\alpha}{2}} \ln \left(1 + \frac{2}{\sqrt{3\alpha}} \right) M_P. \quad (3)$$

In the limit $\alpha \rightarrow \infty$, $\phi_{\text{end}} \rightarrow \sqrt{2} M_P$. The potential at the end of inflation is then given by

$$V_{\text{end}} \equiv V(\phi_{\text{end}}) \simeq \frac{9M^2 M_P^2}{(2\sqrt{3} + 3\sqrt{\alpha})^2}. \quad (4)$$

The mass scale M can be expressed in terms of \mathcal{P}_S (Kallosh et al., 2013; Ozkan et al., 2015)

$$\frac{M}{M_P} = \sqrt{\frac{128\pi^2 \mathcal{P}_S}{3\alpha}} \frac{x_{\text{obs}}}{(1 - x_{\text{obs}})^2}, \quad (5)$$

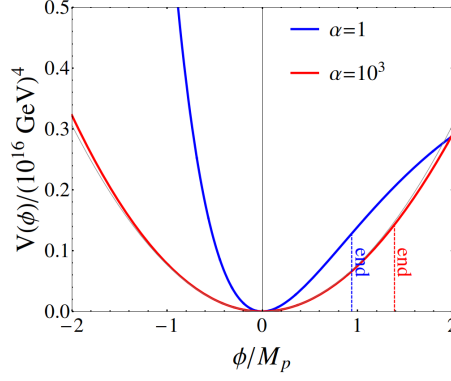


Figure 2: Shape of the potential given by Eq. (2) for $\alpha = 1$ and 10^3 . Here ‘end’ refers to the end of inflation, as given by Eq. (3). The gray curve (below the red curve) denotes a quadratic potential.

where $x_{\text{obs}} \equiv x(\phi_{\text{obs}})$ which can be evaluated once the number of e-foldings N is known (Kallosh et al., 2013; Ozkan et al., 2015):

$$N = \frac{3}{4}\alpha \left(\frac{1}{x_{\text{obs}}} - \frac{1}{x_{\text{end}}} \right) + \left(\frac{3}{4}\alpha - \frac{1}{2} \right) \ln \left(\frac{x_{\text{obs}}}{x_{\text{end}}} \right), \quad (6)$$

where $x_{\text{end}} \equiv x(\phi_{\text{end}}) = (1 + 2/\sqrt{3\alpha})^{-1}$.

Using Eq. (6), the tensor-to-scalar ratio for the potential in Eq. (2) is given by

$$r(\alpha) = \frac{64x_{\text{obs}}^2(\alpha)}{3\alpha[1 - x_{\text{obs}}(\alpha)]^2}. \quad (7)$$

For $\alpha = \mathcal{O}(1)$, $M \simeq \sqrt{24\pi^2\alpha\mathcal{P}_s} M_P/N$ and $r \simeq 12\alpha/N^2$, whereas for $\alpha \gg 1$, $M \simeq \sqrt{6\pi^2\alpha\mathcal{P}_s} M_P/N$ and $r \simeq 8/N$. In Fig. 2 we show a plot of the potential Eq. (2) for $\alpha = 1$ which leads to $r \simeq \mathcal{O}(0.004)$, and $\alpha = 10^3$ which leads to $r \simeq \mathcal{O}(0.1)$. We also show the value of ϕ at which inflation terminates in each case.

In Fig. 3(a), we show the inflaton potential energy at the end of inflation as a function of $r(\alpha)$ for 50 and 60 e-foldings, since the precise value of N depends on the details of reheating process (Lyth and Liddle, 2009). Once inflation ends, the field ϕ starts oscillating around the minimum of the potential with an effective mass

$$m_{\phi,\text{eff}} \simeq \left(\frac{\partial^2 V(\phi)}{\partial \phi^2} \right)^{1/2} \simeq \frac{M}{\sqrt{\alpha}} (2e^{-2z} - e^{-z})^{1/2}, \quad (8)$$

where $z = \sqrt{2/3\alpha} \phi/M_P$. The evolution of ϕ is governed by the following background equation of motion during the inflaton oscillations:

$$\frac{\partial^2 \phi}{\partial t^2} + (3H + \Gamma_\phi) \frac{\partial \phi}{\partial t} + \frac{\partial V(\phi)}{\partial \phi} \simeq 0, \quad (9)$$

where Γ_ϕ is the decay rate of the inflaton field. Eq. (9) holds true as long as $m_\phi \gg H(t)$, so that the thermal and backreaction effects can be neglected. For the potential given by Eq. (2), the amplitude of ϕ oscillation quickly drops with the expansion of the Universe right after the end of inflation; it becomes less than 0.2 of its initial value after only one oscillation. Thus, after a few oscillations, ϕ becomes confined to a small region around the minimum of the potential for which the potential can be approximated by a quadratic one with a constant mass, $m_\phi \equiv m_{\phi,\text{eff}}(\phi \ll M_P) \simeq M/\sqrt{\alpha}$. Note that in the limit $\alpha \rightarrow \infty$, $m_{\phi,\text{eff}}$ globally approaches $M/\sqrt{\alpha} \simeq \sqrt{6\pi^2\mathcal{P}_s} M_P/N$. In Fig. 3(b), we show m_ϕ as a function of r for $N = 50$ and 60.

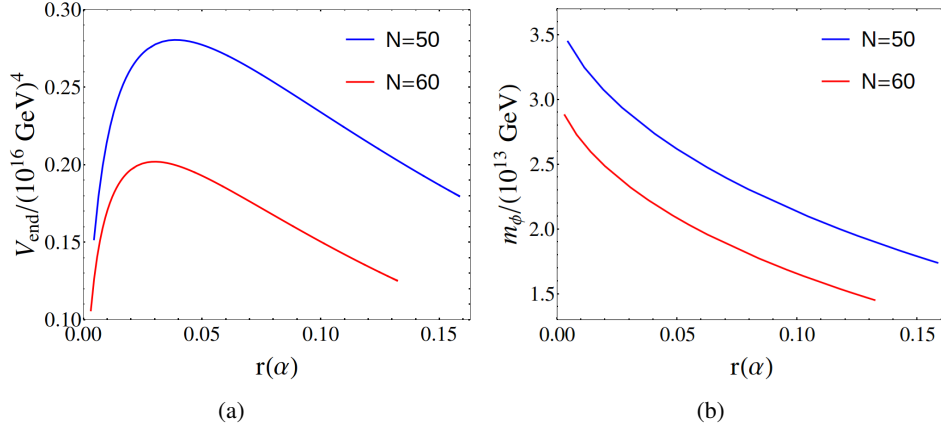


Figure 3: (a) The inflaton energy density at the end of inflation and (b) effective inflaton mass for $\phi \ll M_P$ as a function of $r(\alpha)$ for $N = 50$ (60) shown by the blue (red) curve.

3. DARK MATTER PRODUCTION FROM INFLATON DECAY PRODUCTS

The renormalizable interactions involving the inflaton ϕ , DM χ , heavy mediator X and the relevant SM fields (RH fermions ψ_R and Higgs doublet φ) are given by

$$-\mathcal{L}_{\text{int}} \supset y_{\phi\psi}\phi\bar{\psi}_R\psi_R + \frac{1}{2}y_{\phi\varphi}^2\phi^2\varphi^\dagger\varphi + \frac{1}{2}y_{\phi X}^2\phi^2X^2 + y_{\phi\chi}\phi\bar{\chi}\chi + y_{X\chi}X\bar{\chi}\chi + y_{X\psi}X\bar{\psi}_R\psi_R + \text{H.c.} \quad (10)$$

Note that χ belonging to the hidden sector does not have a direct coupling to the SM fermions and the only indirect coupling arises by integrating out the mediator or the inflaton field in Eq. (10). Therefore, the effective interaction between the SM d.o.f and DM will be determined by a dimension-6, four-Fermion operator $\bar{\chi}\chi\bar{\psi}_R\psi_R/m_\phi^2$ or $\bar{\chi}\chi\bar{\psi}_R\psi_R/m_X^2$, which will be respectively suppressed by the mass square of inflaton or the heavy mediator field, with $\mathcal{O}(1)$ Yukawa couplings. This naturally leads to a FIMP or non-thermal DM scenario which we are interested in here (Dev et al., 2014; Blennow et al., 2014; Baer et al., 2014; Elahi et al., 2015). We will keep our discussion general, without referring to any particular DM model.

In order to have a standard radiation-dominated era just after reheating, we require $y_{\phi\chi}, y_{\phi\varphi}, y_{\phi X} \ll y_{\phi\psi}$. For the sake of simplicity and illustration, we may assume the couplings $y_{\phi\varphi}, y_{\phi X} \approx 0$. Any reasonable value of $y_{\phi\varphi}(y_{\phi X}) \gtrsim 10^{-6}$ would in principle lead to a non-perturbative production of $X(\varphi)$ (Shtanov et al., 1995; Kofman et al., 1994), but in this case inflaton does not decay completely. One would still require the inflaton to decay perturbatively, which will be guaranteed to happen in our case via $y_{\phi\psi}\phi\bar{\psi}_R\psi_R$ (Allahverdi et al., 2010). Thanks to small Yukawa interactions, we can also ignore issues like fermionic preheating (Giudice et al., 1999) or fragmentation of the inflaton (Enqvist et al., 2002b,a). Further, we require that $y_{\phi\psi} \ll 1$ in order to avoid radiative corrections to the inflationary potential (Coleman and Weinberg, 1973) and thermal corrections (Drewes, 2014).

3.1. DM Production

Let us now compute the DM production rates due to both the channels shown in Fig. 1. Besides decaying to DM pairs, the inflaton will decay dominantly to the SM radiation and its total decay rate is given by

$$\begin{aligned} \Gamma_\phi &\simeq \Gamma(\phi \rightarrow \bar{\psi}\psi) + \Gamma(\phi \rightarrow \bar{\chi}\chi) \\ &\simeq \sum_\psi c_\psi y_{\phi\psi}^2 \frac{m_{\phi,\text{eff}}}{8\pi} \left(1 - \frac{4m_{\psi,\text{eff}}^2}{m_{\phi,\text{eff}}^2}\right)^{3/2} + y_{\phi\chi}^2 \frac{m_{\phi,\text{eff}}}{8\pi} \left(1 - \frac{4m_{\chi,\text{eff}}^2}{m_{\phi,\text{eff}}^2}\right)^{3/2}, \end{aligned} \quad (11)$$

where c_ψ denotes the color factor, and $m_{\chi,\text{eff}} \simeq |m_\chi + y_{\phi\chi}\phi|$ and $m_{\psi,\text{eff}} \simeq |m_\psi^{\text{th}} + y_{\phi\psi}\phi|$ are the effective masses of DM and RH SM fermions, respectively, where m_ψ^{th} denotes the plasma induced thermal masses

for the RH SM sector (Weldon, 1982). For $m_{\chi,\text{eff}}, m_{\psi,\text{eff}} \ll m_{\phi,\text{eff}}$ and $\phi \ll M_P$, we can approximate $\Gamma_\phi \simeq \alpha_\phi m_{\phi,\text{eff}} \simeq \alpha_\phi M / \sqrt{\alpha}$. Assuming universal inflaton coupling to RH SM fermions and suppressing the coupling to RH neutrinos (if any), we have $\alpha_\phi \simeq 21 y_{\phi\psi}^2 / 8\pi$, where α_ϕ can be defined in terms of the reheat temperature T_{rh} through the relation (Kolb and Turner, 1990; Chung et al., 1998, 1999)

$$H^2(\tau_\phi) \simeq \frac{\Gamma_\phi^2}{4} \simeq \frac{(\alpha_\phi m_\phi)^2}{4} \simeq \frac{1}{3M_P^2} \frac{\pi^2}{30} g_\rho T_{\text{rh}}^4, \quad (12)$$

which gives the following expression for α_ϕ :

$$\alpha_\phi = \left(\frac{2\pi^2}{45} \right)^{1/2} g_\rho^{1/2} \frac{T_{\text{rh}}^2}{M_P m_\phi}. \quad (13)$$

The inflaton branching ratio to DM is then given by

$$B_\chi \equiv \Gamma(\phi \rightarrow \bar{\chi}\chi) / \Gamma_\phi \sim y_{\phi\chi}^2 / (21 y_{\phi\psi}^2) \ll 1. \quad (14)$$

Note that the finite temperature effects on the inflaton decay rate are negligible as long as the maximum attainable temperature, $T_{\text{max}} \ll m_\phi$, where (Kolb and Turner, 1990; Chung et al., 1999; Mazumdar and Zaldivar, 2014; Drewes, 2014)¹

$$T_{\text{max}} \simeq \frac{1.64}{\sqrt{\pi}} g_\rho^{-1/4} \Gamma_\phi^{1/4} M_P^{1/4} \rho_{\phi,I}^{1/8}, \quad (15)$$

with g_ρ being the number of relativistic d.o.f contributing to the radiation energy density and $\rho_{\phi,I} \simeq V_{\text{end}}$ is the inflaton energy density at the end of inflation. As for the thermal contribution, one can easily show that the relevant contribution comes from the X -mediated DM production in Fig. 1(b). In particular, $\bar{\psi}_R \psi_R \rightarrow \phi \rightarrow \bar{\chi}\chi$ would yield a sub-dominant DM contribution compared to $\bar{\psi}_R \psi_R \rightarrow X \rightarrow \bar{\chi}\chi$. This is due to the fact that both thermal and non-thermal contribution to DM abundance in the inflaton-mediation case are proportional to $y_{\phi\chi}$ [cf. Eq. (14)], which has to be small in order not to non-thermally overproduce DM (Allahverdi and Drees, 2002a; Dev et al., 2014). The thermal contribution due to inflaton mediation dominates only when $T_{\text{max}} \gg m_\phi$, in which case the thermal corrections to inflaton decay rate also become important (Drewes, 2014). On the other hand, such complications do not arise in case of the X -mediation as long as $m_X \gtrsim T_{\text{max}}$.

Assuming $B_\chi \ll 1$, we can trace the time evolution of the inflaton decay products by solving the following set of coupled Boltzmann equations (Kolb and Turner, 1990; Rychkov and Strumia, 2007)²:

$$\frac{d\rho_\phi}{dt} + 3H\rho_\phi = -\Gamma_\phi \rho_\phi \quad (16)$$

$$\frac{d\rho_{\text{rad}}}{dt} + 4H\rho_{\text{rad}} = (1 - B_\chi) \Gamma_\phi \rho_\phi \quad (17)$$

$$\frac{dn_\chi}{dt} + 3Hn_\chi = 2B_\chi \Gamma_\phi \frac{\rho_\phi}{m_\phi} + \sum_\psi \gamma(\bar{\psi}_R \psi_R \rightarrow \bar{\chi}\chi), \quad (18)$$

where ρ_ϕ (ρ_{rad}) denotes the inflaton (radiation) energy density, n_χ is the DM number density, γ is the DM thermal production rate and m_ϕ is the effective inflaton mass averaged over one oscillation, as shown in Fig. 3(b). For simplicity, we focus on the case of heavy mediator i.e. $m_X \gg T_{\text{max}}$. Then on dimensional grounds, the cross section $\sigma(\bar{\psi}_R \psi_R \rightarrow \bar{\chi}\chi) \sim y_{\phi\chi}^2 y_{\psi\chi}^2 T^2 / m_X^4$, and since the DM thermal production rate $\gamma \propto \sigma$, we have

$$\gamma(\bar{\psi}_R \psi_R \rightarrow \bar{\chi}\chi) = \mathcal{J}(\bar{\psi}_R \psi_R \rightarrow \bar{\chi}\chi) \frac{T^8}{m_X^4}, \quad (19)$$

¹In the derivation of Eq. (15), the masses of daughter particles are assumed to be much smaller than that of inflaton.

²This set of Boltzmann equations is valid only when ϕ is oscillating around a quadratic minimum, which is quickly realized in our case only after a few oscillations, when $\phi \ll M_P$.

where $\mathcal{J} \sim y_{\chi\chi}^2 y_{\chi\psi}^2$ is a constant for $T \gg m_{\chi,\text{eff}}, m_{\psi,\text{eff}}$ and the factor T^6 arises from the two ψ_R 's being in a thermal plasma. In the numerical calculations we use the exact integral expression for γ (Gondolo and Gelmini, 1991) and also take into account the plasma induced thermal masses for RH SM fermions (Weldon, 1982).

4. THERMAL AND NON-THERMAL DARK MATTER ABUNDANCE

The total DM relic abundance for temperatures $T \ll T_{\text{rh}}$ is roughly given by the sum of the thermal (th) and non-thermal (non-th) components (Gelmini and Gondolo, 2006; Gelmini et al., 2006):

$$\Omega_\chi h^2 \simeq \Omega_\chi^{\text{non-th}} h^2 + \Omega_\chi^{\text{th}} h^2 \simeq 2.74 \times 10^8 (n_\chi^{\text{non-th}} + n_\chi^{\text{th}}) \frac{m_\chi}{s}, \quad (20)$$

where h is the scaled Hubble rate, $s = (2\pi^2/45)g_s T^3$ is the entropy density with g_s being the corresponding number of relativistic d.o.f. We take $g_\rho = g_s \equiv g$, which is valid for most of the thermal history of the Universe. Eqs. (16)-(18) can be simplified by defining comoving energy and number densities (Chung et al., 1999): $\Phi \equiv \rho_\phi a^3$, $R \equiv \rho_{\text{rad}} a^4$ and $X \equiv n_\chi a^3$ where a is the scale factor. The temperature of the plasma during inflaton oscillations reaches its maximum value T_{max} given by Eq. (15) at roughly $a_{\text{max}}/a_I \sim 1.5$, where a_I denotes the initial scale factor at the end of inflation. Thereafter, the temperature decreases with the expansion as $a^{-3/8}$ until radiation takes over signalling the end of reheating phase at a_{rh} . Similarly, the DM thermal and non-thermal abundances, $Y_\chi^{\text{th(non-th)}} \equiv n_\chi^{\text{th(non-th)}}/s$, increase initially very fast to reach their maxima very soon after the beginning of the reheating process and then decrease as $a^{-3/8}$ till the end of reheating.

In the non-thermal case [Fig. 1(a)], the DM relic abundance is directly determined from the inflaton energy density, which it turn depends on $r(\alpha)$ [see the first term on the RHS of Eq. (18) and Fig. 3(a)], besides the sensitivity of the decay rate to the shape of the inflationary potential around the minimum which also depends on $r(\alpha)$. Thus, the connection between the DM properties and the primordial fluctuations is straightforward, but to our knowledge, this important point was never discussed in the literature. We further show that such an interesting connection also exists in case (b) [see Fig. 1(b)] if the mediator mass is between T_{max} and M_P . This can be understood from Eqs. (18) and (19), where the term sourcing the DM thermal abundance is proportional to $T^8 \sim \rho_{\text{rad}}^2$ and since the thermal bath itself arises from the decay of inflaton [cf. Eq. (17)], a connection between DM thermal abundance and the tensor-to-scalar ratio can be established.

In the Appendix, we have obtained approximate analytical expressions for the DM thermal and non-thermal abundances, and argued that in both cases, the details of the inflationary potential are carried over to DM abundance via both m_ϕ and V_{end} , thereby establishing a connection between the DM abundance and tensor-to-scalar ratio. In order to precisely capture this connection between r and $\Omega_\chi h^2$, we need to integrate Eqs. (16)-(18) numerically for different $r(\alpha)$ values. The resulting DM thermal and non-thermal abundances are shown in Figs. 4 (a) and (b), respectively, for a typical choice of parameters: $m_\chi = 0.5$ GeV, $y_{\chi\chi} = y_{\chi\psi} = 1$ and $y_{\phi\psi} = 3.6 \times 10^{-7}$. We find that for the thermal case, $\Omega_\chi^{\text{th}} \propto r^{-3/8}$, while for the non-thermal case, $\Omega_\chi^{\text{non-th}} \propto r^{4/5}$. This is a very interesting connection, especially the former one, since it relates thermal production of DM at the time of reheating and thermalization of the Universe with the value of r which, as a matter of fact, provides a promising probe for the FIMP-like scenario in near future. In other words, if the tensor-to-scalar ratio is measured in future, one can test the properties of FIMP DM through this connection in the context of a given inflationary model.

To examine the dependence on the mediator mass, we scan the (m_χ, m_X) parameter space for the correct DM abundance while fixing the mediator couplings to DM and the SM d.o.f. In Fig. 5, we show the DM thermal abundance heat map as a function of m_χ and m_X for $r \simeq 0.004$. We have fixed $y_{\phi\psi} = 3.6 \times 10^{-7}$, so that $T_{\text{rh}} \simeq 10^9$ GeV. For larger r , the allowed range of (m_χ, m_X) shifts to larger values of m_χ and smaller m_X values. Similarly, for smaller $y_{\phi\psi}$, i.e. smaller reheating temperature, the allowed region of the parameter space shifts to larger DM masses and smaller mediator masses, and vice-versa. For very small values of the branching fraction $B_\chi \ll 1$, the thermal contribution given by Eq. (24) can be dominant over the non-thermal contribution given by Eq. (23), and can account for the observed DM

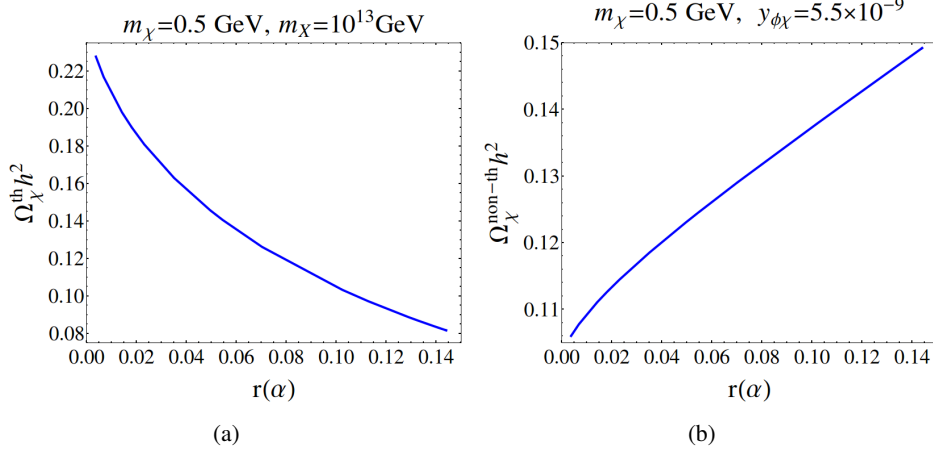


Figure 4: Dark matter (a) thermal and (b) non-thermal abundance as a function of the tensor-to-scalar ratio. Here we fix $m_\chi = 0.5$ GeV, $y_{\chi\chi} = y_{\chi\psi} = 1$ and $y_{\phi\psi} = 3.6 \times 10^{-7}$.

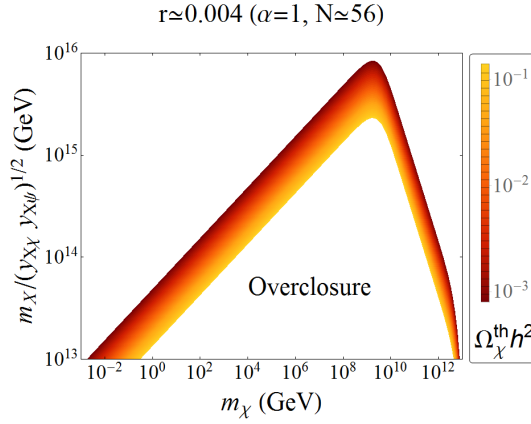


Figure 5: Map of the DM thermal abundance as a function of m_χ and m_X for $r \simeq 0.004$ and $T_{\text{th}} \simeq 10^9$ GeV.

abundance for m_χ as low as roughly $300 \text{ MeV}/(y_{\chi\chi} y_{\chi\psi})^2$. The unshaded region labelled by overclosure (below the colored region) gives $\Omega_\chi h^2 > 0.13$, which is ruled out at 3σ by the latest Planck data (Ade et al., 2015b). The rest of the unshaded region (above the colored region) is still allowed, though for practical purposes, the corresponding DM thermal abundance becomes negligible, and one has to allow for a non-thermal contribution to the DM abundance or to invoke a multi-component DM to explain the observed abundance, see e.g. (Chialva et al., 2013). Note that for $m_\chi \gtrsim T_{\text{th}}$, the thermal production rate is suppressed due to a smaller phase space. This can be seen from the right parts of the allowed (m_χ, m_X) parameter space in Fig. 5.

5. DISCUSSION AND CONCLUSION

The BICEP2 collaboration had recently claimed a positive detection of the primordial tensor modes with more than 5σ significance (Ade et al., 2014). However, a more careful analysis of the foreground dust weakened the claim to an upper limit of $r < 0.12$ at 95% CL (Ade et al., 2015a). Nevertheless, if the tensor-to-scalar ratio is definitively measured in future, it will provide an important indirect handle on the dark sector physics, as shown above, along with a clear evidence of a quantum nature of gravity (Ashoorioon et al., 2014).

Before we conclude, let us briefly mention that we could have also imagined a similar mechanism for producing baryon/lepton (B/L) asymmetry, either from the direct decay of the inflaton, if the inflaton were carrying any B/L number (Murayama et al., 1993), or from the intermediate condensate X carrying

B/L number (Enqvist and Mazumdar, 2003), as e.g. in GUT-baryogenesis (Kolb and Turner, 1990). As a concrete example, in non-supersymmetric models, we can realize high-scale thermal leptogenesis via the production and decay of RH heavy Majorana neutrinos (Davidson et al., 2008). In either case, we would be able to relate the scale of inflation, and therefore the tensor-to-scalar ratio, with the magnitude of the L-asymmetry. In the latter case, one would require a weak washout regime in order to retain the sensitivity towards the initial conditions. A detailed discussion of these issues will be given elsewhere.

To conclude, if the future observations could pin down the exact value of the tensor-to-scalar ratio, it would serve as an interesting way to constrain the hidden sector, including the properties of the DM feebly interacting with the SM d.o.f, which are otherwise very hard to probe. Although for the sake of illustration we have used a particular class of inflationary potential to derive our results, the idea of connecting the DM abundance to the primordial tensor perturbations should hold true for a generic model of inflation, including multi field driven inflationary models (Liddle et al., 1998).

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APPENDIX: ANALYTICAL ESTIMATION OF THE DM ABUNDANCE

With the assumptions mentioned in Section 3.1, namely, (i) $m_\phi \ll M_P$, (ii) $y_{\phi\psi}, y_{\phi\chi} \ll m_\phi/M_P$, and (iii) $m_\psi^{\text{th}}, m_\chi \ll m_\phi$, we derive analytical expressions for the DM abundance produced either thermally or non-thermally.

Non-thermal abundance: Deep inside the inflaton domination (id) epoch, i.e when $H \gg \Gamma_\phi$, $\Phi \simeq \Phi_I$ where $\Phi_I = \rho_{\phi,I} a_I^3 \simeq V_{\text{end}} a_I^3$ and $R_I \simeq X_I \simeq 0$; so Eq. (17) can be easily solved to yield

$$\rho_{\text{rad}}^{\text{id}}(a) \simeq \frac{2\sqrt{3}}{5} \Gamma_\phi M_P V_{\text{end}}^{1/2} \left[\left(\frac{a_I}{a} \right)^{3/2} - \left(\frac{a_I}{a} \right)^4 \right]. \quad (21)$$

This allows us to estimate the temperature of the ambient relativistic d.o.f during the inflaton domination epoch (Chung et al., 1999):

$$T^{\text{id}}(a) = \left(\frac{30}{\pi^2 g_\rho} \rho_{\text{rad}}^{\text{id}} \right)^{1/4} \simeq \left(\frac{432}{\pi^4 g_\rho^2} \right)^{1/8} \Gamma_\phi^{1/4} M_P^{1/4} V_{\text{end}}^{1/8} \left[\left(\frac{a_I}{a} \right)^{3/2} - \left(\frac{a_I}{a} \right)^4 \right]^{1/4}. \quad (22)$$

Now Eq. (18) can be easily solved for the DM non-thermal component, which is sourced by the first term on the RHS of Eq. (18). Evaluating this general expression for the non-thermal number density at T_{rh} , and accounting for the inflaton population decaying at $T < T_{\text{rh}}$ and the accompanying entropy release (Giudice et al., 2001), we reproduce the non-thermal contribution to the relic density (Dev et al., 2014; Allahverdi and Drees, 2002a)

$$\begin{aligned} \frac{\Omega_\chi^{\text{non-th}} h^2}{0.12} &\simeq 3.87 \times 10^5 B_\chi \left(\frac{g}{106.75} \right)^{-1/4} \left(\frac{\alpha_\phi}{10^{-13}} \right)^{1/2} \left(\frac{m_\chi}{1 \text{ GeV}} \right) \left(\frac{m_\phi}{10^{13} \text{ GeV}} \right)^{-3/4} \left(\frac{V_{\text{end}}}{(10^{16} \text{ GeV})^4} \right)^{1/8} \\ &\simeq 4.44 \times 10^5 B_\chi \left(\frac{m_\chi}{1 \text{ GeV}} \right) \left(\frac{m_\phi}{10^{13} \text{ GeV}} \right)^{-1} \left(\frac{T_{\text{rh}}}{10^9 \text{ GeV}} \right). \end{aligned} \quad (23)$$

Here we have used Eq. (13) for α_ϕ and Eq (22) to change the dependence from a_{rh} to T_{rh} . Clearly, the DM non-thermal abundance depends on the inflaton energy at the end of inflation, V_{end} , and it is also sensitive to the steepness of the inflationary potential around the minimum characterized by m_ϕ . In general, both V_{end} and m_ϕ depend on $r(\alpha)$ [cf. Fig. 3] for the class of potentials under consideration. This gives rise to the $\Omega_\chi^{\text{non-th}} \propto r^{4/5}$ behavior in Fig. 4(b).

Thermal contribution: For thermal contribution, since $\gamma \propto T^8$ [cf. Eq. (19)], which during the radiation domination epoch will redshift as a^{-8} , the scaled version of Eq. (18) will in turn go as $d[n_\chi^{\text{th}} a^3]/da \propto a^{-4}$. This means that except for a dilution factor due to the entropy released directly after the transition to the radiation domination phase, the DM thermal yield becomes constant directly after the end of reheating $Y_\chi^{\text{th}}(T \ll T_{\text{rh}}) \simeq \zeta Y_\chi^{\text{th}}(T = T_{\text{rh}})$, where ζ is the dilution factor. Thus, we obtain

$$\begin{aligned} \frac{\Omega_\chi^{\text{th}} h^2}{0.12} &\simeq 1.73 y_{\chi\chi}^2 y_{\chi\psi}^2 \left(\frac{g}{106.75}\right)^{-9/4} \left(\frac{\alpha_\phi}{10^{-13}}\right)^{3/2} \left(\frac{m_\chi}{1 \text{ GeV}}\right) \left(\frac{m_X}{10^{13} \text{ GeV}}\right)^{-4} \left(\frac{m_{\phi,\text{eff}}}{10^{13} \text{ GeV}}\right)^{3/4} \left(\frac{V_{\text{end}}}{(10^{16} \text{ GeV})^4}\right)^{3/8} \\ &\simeq 2.61 y_{\chi\chi}^2 y_{\chi\psi}^2 \left(\frac{g}{106.75}\right)^{-3/2} \left(\frac{m_\chi}{1 \text{ GeV}}\right) \left(\frac{m_X}{10^{13} \text{ GeV}}\right)^{-4} \left(\frac{T_{\text{rh}}}{10^9 \text{ GeV}}\right)^3. \end{aligned} \quad (24)$$

Here again we have used Eq. (13) for α_ϕ and Eq. (22) to change the dependence from a_{rh} to T_{rh} . It is clear that the DM thermal abundance is also sensitive to the inflaton energy at the end of inflation and inflaton mass around the minimum of the potential, see Figs. 3(a) and 3(b). Hence follows the $\Omega_\chi^{\text{th}} \propto r^{-3/8}$ behavior in Fig. 4(a).

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